

# Penetration and amplification of resonant perturbations in 3D ideal-MHD equilibria

S. R. Hudson<sup>1</sup>, J. Loizu<sup>2,1</sup>, S. Lazerson<sup>1,2</sup>, A. Bhattacharjee<sup>1</sup>, and P. Helander<sup>2</sup>

<sup>1</sup> Princeton Plasma Physics Laboratory, PO Box 451, Princeton NJ 08543, USA

<sup>2</sup> Max-Planck-Institut für Plasmaphysik, D-17491 Greifswald, Germany

(Dated: September 26, 2016)

The nature of ideal MHD equilibria in three-dimensional geometry is profoundly affected by resonant surfaces, and this is particularly true for equilibria with non-zero pressure. If not treated carefully, infinite, non-physical currents arise in equilibria with continuously-nested magnetic surfaces. We demonstrate that three-dimensional, ideal-MHD equilibria, with nested surfaces, and with  $\delta$ -function current-densities that produce a *discontinuous* rotational-transform, are well defined and can be computed both perturbatively and using fully-nonlinear equilibrium calculations. The results are of direct practical importance: we predict that resonant magnetic perturbations penetrate past the rational surface (i.e. “shielding” is incomplete, even in purely ideal-MHD) and that the perturbation is amplified by plasma pressure, increasingly so as stability limits are approached.

## I. IMPORTANCE OF $\nabla p = \mathbf{j} \times \mathbf{B}$

The properties and numerical computation of three-dimensional (3D), ideal-MHD equilibria, as described by the seemingly-simple yet deceptively-complicated, ideal force-balance equation

$$\nabla p = \mathbf{j} \times \mathbf{B}, \quad (1)$$

is of fundamental importance for understanding the behavior of both magnetically-confined fusion and astrophysical plasmas [1, 2]. The potentially-dangerous instabilities that cause edge-localized modes (ELMs), an important concern for ITER [3] are widely believed to be ideal, peeling-ballooning modes [4]; and a ‘hot-topic’ of current research is to discover how these modes may be suppressed by applying resonant magnetic perturbations (RMPs), i.e. by 3D effects [5]. Whilst there exist extended-MHD codes [6] with non-ideal physics, these codes come with significant computational cost; and so the plasma response to 3D perturbations is routinely determined using codes such as IPEC [7], the ideal, perturbed equilibrium code.

However, there are two fundamental difficulties that are frequently over-looked in calculations of ideal-MHD equilibria with nested flux-surfaces and smooth pressure and rotational-transform profiles. The first is the existence of *infinite* currents near resonant, rational rotational-transform surfaces, and the second is that ideal-MHD equilibria are *not* analytic functions of the 3D boundary.

Our new class of solutions for ideal-MHD equilibria avoids these problems. Most importantly, our solutions yield predictions that are in sharp contrast to previous predictions, and which may have direct implications for understanding the effect of RMPs in tokamaks. Even in ideal-MHD, a resonant perturbation penetrates past the rational surface and into the core of the plasma; and the perturbation is magnified by pressure inside the resonant surface, increasingly so as stability limits are approached.

The content and outline of this paper is as follows. As our new class of solutions and their properties have already been described in detail in earlier publications [8, 9], this article shall review the salient features of these new solutions from an intuitive perspective and raise (and hopefully clarify) some of the problematic issues that remain. To appreciate the new class of solu-

tions, it is required to understand the difficulties with the conventional solutions; so, this paper first provides a review of the key problem of equilibria with nested flux surfaces and smooth pressure gradients, namely the formation of *unphysical* currents near rational surfaces. Then, we briefly review the problem of *non-analyticity* of ideal-MHD, and how this undermines perturbative approaches; and we provide an illustration in cylindrical geometry.

## II. UNPHYSICAL, INFINITE CURRENTS

The infinite, unphysical currents arise from enforcing charge conservation,  $\nabla \cdot \mathbf{j} = 0$ , with smooth pressure profiles. It is convenient to write the current density as  $\mathbf{j} \equiv u\mathbf{B} + \mathbf{j}_\perp$ , and by crossing  $\nabla p = \mathbf{j} \times \mathbf{B}$  with  $\mathbf{B}$  we derive an expression for the perpendicular current density,

$$\mathbf{j}_\perp = \frac{\mathbf{B} \times \nabla p}{B^2}. \quad (2)$$

Wherever there are pressure-driven, perpendicular current-densities, there must also be parallel current-densities that satisfy  $\mathbf{B} \cdot \nabla u = -\nabla \cdot \mathbf{j}_\perp$ . Such equations, called magnetic differential equations [10], take a simple form in straight-field line coordinates,  $(\psi, \theta, \zeta)$ , for which

$$\mathbf{B} = \nabla\psi \times \nabla\theta + \iota(\psi)\nabla\zeta \times \nabla\psi, \quad (3)$$

where  $2\pi\psi$  is the toroidal flux enclosed by a flux surface,  $\theta$  and  $\zeta$  are straight fieldline poloidal and toroidal angles, and  $\iota(\psi)$  is the rotational-transform. (Note that  $\mathbf{B} \cdot \nabla\zeta = \sqrt{g}^{-1}$ , where  $\sqrt{g}$  is the coordinate Jacobian.) This is a linear equation, and may be solved in Fourier space by the representation  $u = \sum u_{m,n} \exp(im\theta - in\zeta)$ . The directional-derivative  $\sqrt{g}\mathbf{B} \cdot \nabla = \iota\partial_\theta + \partial_\zeta$ , reduces to  $\sqrt{g}\mathbf{B} \cdot \nabla = i(\iota m - n)$ . The solution for each Fourier harmonic of the parallel current is

$$u_{m,n} = \frac{i(\sqrt{g}\nabla \cdot \mathbf{j}_\perp)_{m,n}}{x} + \Delta_{m,n} \delta(x), \quad (4)$$

where  $x \equiv \iota m - n$  is the “distance” in rotational-transform from the rational surface, and  $\Delta_{m,n}$  is an as-yet undetermined constant.

Two singularities are present: a pressure-driven, so-called “Pfirsch-Schlüter”,  $1/x$  current-density that arises

around rational surfaces, and a  $\delta$ -function current-density that develops at rational surfaces.

Singularities in the current-*density* are acceptable. As a simple example, consider a finite, non-zero current passing along a thin wire: if the cross-sectional area of the wire is very thin, then representing the current density as a  $\delta$ -function is an excellent mathematical approximation. In the limit that the conductivity becomes infinite, a wire of zero cross-sectional area can support a finite current, and the  $\delta$ -function approximation becomes exact. The existence of  $\delta$ -function currents, which are also sometimes called “sheet” currents or “surface” currents, is perfectly compatible with the physical model of ideal-MHD.

Singularities in the *current*, however, are not acceptable: the total current,  $\int \mathbf{j} \cdot d\mathbf{s}$ , passing through *each* and *every* surface must remain finite for any physically acceptable equilibrium. The current resulting from the pressure-driven Pfirsch-Schlüter density passing through the cross-sectional area enclosed by the flux-surfaces  $x = \epsilon$  and  $x = \delta$  between  $\theta = 0$  and  $\theta = \pi/m$  is proportional to  $\int_{\epsilon}^{\delta} 1/x dx$ , which logarithmically approaches infinity as  $\epsilon$  approaches zero.

This is not physical. The presence of an equal-and-opposite infinite current, in a sense, through a similar cross-sectional surface on the opposing side of the rational surface (i.e., the surface bounded by  $x = -\epsilon$  and  $x = -\delta$ ) should not be thought of as a panacea. It would seem that the ideal-MHD equilibrium model, with nested flux-surfaces, has the fatal flaw of not allowing for pressure gradients in a small-but-finite neighborhood of each rational surface! Considering that rational surfaces are dense in any continuous magnetic field with shear, and therefore the regions in which the pressure gradient must be zero will overlap, this means that there can no pressure at all in equilibria with nested flux surfaces and smooth pressure (unless the rotational-transform is everywhere irrational, as we discuss in more detail below).

There are other equilibrium models that should be mentioned, but which will not be discussed at length in this article. If the pressure gradient is allowed to be *discontinuous*, a pressure profile with finite gradients where the rotational-transform is sufficiently irrational will provide a non-trivial, finite-pressure equilibrium. For example: choose  $p'(\psi) = 1$  where the rotational-transform satisfies a Diophantine condition, e.g.  $|\iota(\psi) - n/m| > k/m^2$  for all rationals  $n/m$  and some  $k > 0$ , and  $p'(\psi) = 0$  elsewhere; such a profile would avoid the pressure-driven infinite-currents near the rational surfaces, yet would yield a non-trivial pressure profile (by virtue of the fact that a finite measure of irrationals satisfy the Diophantine condition). In this case, however, the equilibrium must display “pathological” fractal properties, [1], and would therefore not be amenable to standard numerical discretization (it is almost impossible to constrain the topology of a non-integrable field to be consistent with a given, *fractal* pressure). Other possibilities include equilibria in which the pressure is discontinuous, as is assumed in the Stepped Pressure Equilibrium Code (SPEC) [11], and equilibria for which the boundary is appropriately constrained to ensure that the resonant harmonics of the geometry vanish in the vicinity of the rational surfaces [12–14].

### III. NON-ANALYTICITY

The second problem, the non-analyticity of ideal-MHD equilibria, is encountered when computing the plasma displacement,  $\boldsymbol{\xi} = \xi^{\psi} \mathbf{e}_{\psi} + \xi^{\theta} \mathbf{e}_{\theta} + \xi^{\zeta} \mathbf{e}_{\zeta}$ , induced by an  $\mathcal{O}(\epsilon)$  perturbation to the boundary of an equilibrium state. Assuming an expansion

$$\boldsymbol{\xi} \equiv \epsilon \boldsymbol{\xi}_1 + \epsilon^2 \boldsymbol{\xi}_2 + \dots, \quad (5)$$

the first-order term satisfies  $L_0[\boldsymbol{\xi}_1] = 0$ , where

$$L_0[\boldsymbol{\xi}] \equiv \nabla \delta p - \delta \mathbf{j} \times \mathbf{B} - \mathbf{j} \times \delta \mathbf{B}, \quad (6)$$

where the ideal variations in the magnetic field and pressure are  $\delta \mathbf{B}[\boldsymbol{\xi}] \equiv \nabla \times (\boldsymbol{\xi} \times \mathbf{B})$  and  $\delta p[\boldsymbol{\xi}] = (\gamma - 1) \boldsymbol{\xi} \cdot \nabla p - \gamma \nabla \cdot p \boldsymbol{\xi}$ .

The linear operator,  $L_0[\boldsymbol{\xi}]$ , is singular. To simultaneously match a non-trivial condition at the boundary and the condition that  $\delta \mathbf{B} \cdot \nabla \psi = 0$  at the rational surface (so that a magnetic island does not form), the solution for  $\boldsymbol{\xi}_1$  must, generally, be discontinuous [8].

A discontinuous plasma displacement is, however, inconsistent with the assumption of nested flux-surfaces: in fact, magnetic surfaces overlap if the displacement anywhere violates  $|d\xi/d\psi| < 1$ . The second order term satisfies  $L_0[\boldsymbol{\xi}_2] = -\delta \mathbf{j}[\boldsymbol{\xi}_1] \times \delta \mathbf{B}[\boldsymbol{\xi}_1]$ , and  $\boldsymbol{\xi}_2$  is even more singular than  $\boldsymbol{\xi}_1$ . That perturbation theory is not a valid approach for computing ideal equilibria with nested surfaces was known by Rosenbluth *et al.* [15], who wrote “we must abandon the perturbation theory approach” when computing ideal-MHD equilibria in 3D.

These problems are partly associated with what might be considered fundamental flaws in ideal-MHD [16], namely that ideal-MHD dynamical evolutions do not allow the topology of the magnetic field to “tear”. Despite such flaws in ideal-MHD, the equation of force balance,  $\nabla p = \mathbf{j} \times \mathbf{B}$ , which this paper is primarily concerned with, remains an excellent description of the macroscopic forces that determine plasma equilibrium properties. Considering first the presumably simpler task, that of only seeking *static* solutions, and ignoring the more-complicated questions regarding ideal *dynamical* evolution, the main obstruction to computing mathematically self-consistent solutions to  $\nabla p = \mathbf{j} \times \mathbf{B}$  with nested flux surfaces and smooth pressure profiles was that well-defined solutions had not, until quite recently, been discovered. There is, surprisingly, a class of ideal-MHD equilibria with nested magnetic surfaces and arbitrary pressure profiles that eliminates both the infinite currents and the non-analyticity.

In recent work [17] we, for the first time, computed the  $1/x$  and  $\delta$ -function singular current-densities in 3D equilibria; and we realized that self-consistent solutions demand locally-infinite shear at the resonant surfaces. We then [8] introduced a new class of solution that admits additional  $\delta$ -function current-densities that produce finite sheet currents, with a commensurate *discontinuity* in the rotational-transform.

#### A. Illustration in cylindrical geometry

For illustration and verification, we consider the linear, ideal response to a resonant boundary perturbation in

cylindrical geometry, with “major” radius  $R$  and minor radius  $a$ . (Note, as mentioned above, we do not explicitly follow the ideal dynamical response, but instead simply solve  $\nabla p = \mathbf{j} \times \mathbf{B}$  with the perturbed boundary, i.e. we seek the static solution.) The cylindrically symmetric solution to  $\nabla p = \mathbf{j} \times \mathbf{B}$  satisfies

$$\frac{dp}{dr} + \frac{1}{2} \frac{d}{dr} \left[ B_z^2 \left( 1 + \iota^2 \frac{r^2}{R^2} \right) \right] + \frac{r \iota^2 B_z^2}{R^2} = 0, \quad (7)$$

where  $\iota \equiv RB_\theta/rB_z$ , and is uniquely determined by the value of the axial field at the origin, e.g.  $B_z(0) = 1$ , by an arbitrary pressure profile,  $p(r)$ , and a rotational-transform profile,  $\iota(r)$ . This equation may be cast as an ordinary differential equation for  $B_z$ , and can be integrated from  $r = 0$  to  $r = a$ .

In cylindrical geometry there is no coupling, to lowest order, between incommensurate perturbations; and we may for simplicity of illustration imagine that there is a single resonant surface of concern. (In a fully three-dimensional configuration, all rational surfaces within the equilibrium will potentially create problems.) We therefore choose a rotational-transform profile

$$\iota(r) = \begin{cases} \iota_0 - \iota_1(r/a)^2 + \Delta\iota/2, & \text{for } r < r_s, \\ \iota_0 - \iota_1(r/a)^2 - \Delta\iota/2, & \text{for } r > r_s, \end{cases} \quad (8)$$

with  $\iota_0$  and  $\iota_1$  chosen so that  $\iota(r)$  jumps across the rational  $\iota_s \equiv n/m$ , namely  $\iota(r_s) = \iota_s \pm \Delta\iota/2$ . The pressure can be an arbitrary smooth function, e.g.

$$p(r) = p_0[1 - 2(r/a)^2 + (r/a)^4], \quad (9)$$

and for later reference we define  $\beta \equiv 2p(0)/B_z^2(0)$ .

In cylindrical geometry, the linearized equation,  $L_0[\xi] = 0$ , reduces to Newcomb’s equation [18],

$$\frac{d}{dr} \left( f \frac{d\xi}{dr} \right) - g\xi = 0, \quad (10)$$

where  $\xi^r \equiv \xi(r) \cos(m\theta - n\zeta)$ , and  $f$  and  $g$  are determined by the equilibrium [8]. For  $\Delta\iota = 0$ , Newcomb’s equation is singular where  $\iota(r_s) = n/m$ , because there is a order-two zero in  $f(r)$  and an order-one zero in  $g(r)$ . For  $\Delta\iota \neq 0$ , the singularity is removed; and the equilibrium equations comprise an analytic function of the 3D boundary and the perturbation expansion is valid.

Fig.1 shows the result of numerical integration of Eq. (10), with different values of  $p_0$  and for  $\Delta\iota = 10^{-3}$ . The linear radial displacement is continuous and smooth provided  $\Delta\iota \neq 0$ , and the condition  $|d\xi/d\psi| < 1$  is satisfied if  $\Delta\iota > \Delta\iota_{min}$ , which can be estimated theoretically [8]. Even for a small, local change in the transform profile the global solution is significantly different and the displacement penetrates inside the resonant surface and into the origin. In fact, for non-zero pressure, the perturbation to the displacement inside the resonant surface is magnified, extremely so as ideal-stability limits are approached. To quantify the amplification and penetration of the perturbation, we define

$$A_{rmp} \equiv \xi_s/\xi_a, \quad (11)$$

$$P_{rmp} \equiv 1 - r_*/r_s, \quad (12)$$

where  $r_*$  is the radius at which  $\xi(R_*) = \xi_s/e$ . In Fig.2 and Fig.3 we show these quantities as a function of  $\beta$ . These figures show that as the equilibrium approaches the Suydam limit for interchange stability, given by

$$D_S = \left( \frac{2p'\iota^2}{rB_z^2\iota'^2} \right)_s < \frac{1}{4}, \quad (13)$$

there is a fantastic increase in both the amplification and penetration of the resonant magnetic perturbation.

## B. Verification of a nonlinear equilibrium code

Removing the singularity associated with the rational surface (achieved by allowing for a sheet current that results in a discontinuous rotational-transform) means that ideal-MHD equilibria become analytic functions of the three-dimensional boundary, and this in turn allows nonlinear 3D ideal-MHD equilibrium codes to be verified against the linear codes. Presently, the widely-used, 3D, nonlinear ideal-MHD equilibrium codes VMEC [19] and NSTAB [20] are restricted to work with smooth functions and cannot, formally, compute equilibria with discontinuous rotational-transform (though, finite radial resolution seems to imply an ‘effective’ discontinuity [21]). The SPEC [11] code *does* allow for discontinuities. In Fig.4, we compare the linear plasma displacement computed using SPEC to analytic solution, with excellent agreement. Furthermore, fully *nonlinear* SPEC solutions were compared to the linear solutions [8], obtaining the as-expected agreement.

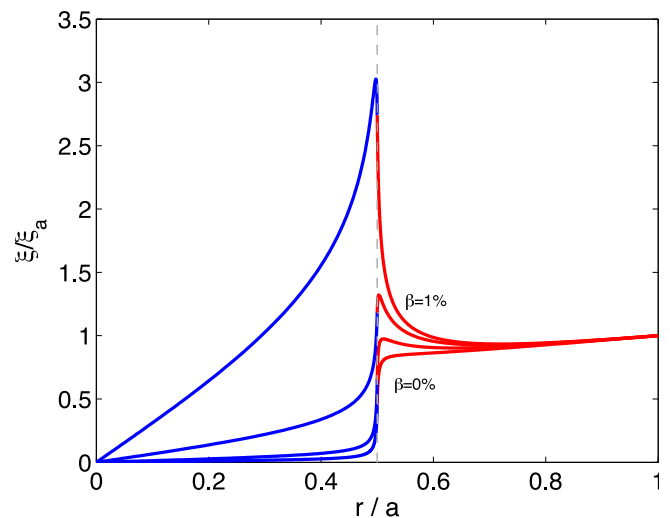


FIG. 1: Solutions for the perturbed displacement for  $\beta = 0\%$  (lower curve) to  $\beta = 1\%$  (upper curve); and for  $\Delta\iota = 10^{-3}$ . (Reproduced from [9].)

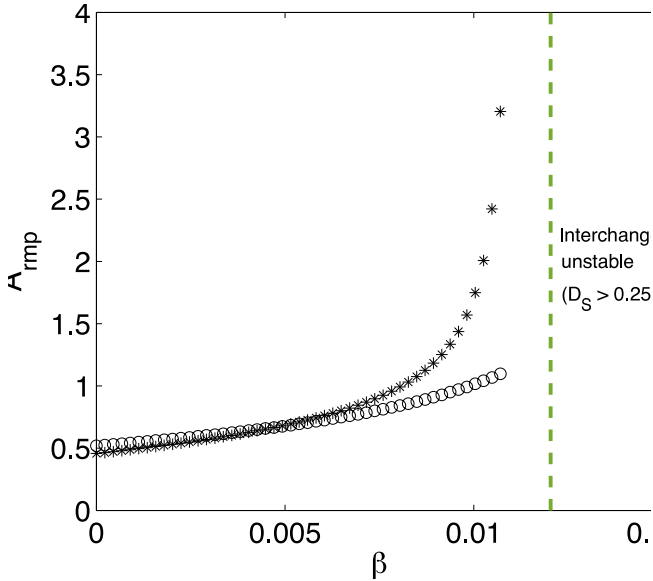


FIG. 2: Amplification of the perturbation on the resonant surface as a function of  $\beta$ , for  $\Delta t = 0.005$  (circles) and  $\Delta t = 0.001$  (stars). The vertical dashed line indicates where the Suydam condition is not satisfied. (Reproduced from [9].)

- [1] H. Grad. Toroidal containment of a plasma. *Phys. Fluids*, 10(1):137, 1967.
- [2] E. Parker. Spontaneous current sheets in magnetic fields. *Oxford University Press*, 1994.
- [3] K. Tomabechi, J. R. Gilleland, Yu. A. Sokolov, R. Toschi, and the ITER team. Iter conceptual design. *Nucl. Fus.*, 31(6):1135, 1991.
- [4] P.B. Snyder, H.R. Wilson, J.R. Ferron, L.L. Lao, A.W. Leonard, T.H. Osborne, A.D. Turnbull, D. Mossessian, M. Murakami, and X.Q. Xu. Edge localized modes and the pedestal: A model based on coupled peeling-ballooning modes. *Phys. Plasmas*, 9(5):2037, 2002.
- [5] T. E. Evans, R. A. Moyer, P. R. Thomas, J. G. Watkins, T. H. Osborne, J. A. Boedo, E. J. Doyle, M. E. Fenstermacher, K. H. Finken, R. J. Groebner, M. Groth, J. H. Harris, R. J. La Haye, C. J. Lasnier, S. Masuzaki, N. Ohya, D. G. Pretty, T. L. Rhodes, H. Reimerdes, D. L. Rudakov, M. J. Schaffer, G. Wang, and L. Zeng. Suppression of large edge-localized modes in high-confinement DIII-D plasmas with a stochastic magnetic boundary. *Phys. Rev. Lett.*, 92(23):235003–1, 2004.
- [6] S. C. Jardin, J. Breslau, and N. Ferraro. A high-order implicit finite element method for integrating the two-fluid magnetohydrodynamic equations in two dimensions. *J. Comp. Phys.*, 226:2146, 2007.
- [7] J.-K. Park, A. H. Boozer, and A. H. Glasser. Computation of three-dimensional tokamak and spherical torus equilibria. *Phys. Plasmas*, 14(5):052110, 2007.
- [8] J. Loizu, S. R. Hudson, A. Bhattacharjee, S. Lazerson, and P. Helander. Existence of three-dimensional ideal-mhd equilibria with current sheets. *Phys. Plasmas*, 22:090704, 2015.
- [9] J. Loizu, S. R. Hudson, Helander, Lazerson, and Bhattacharjee. Pressure-driven amplification and penetration of resonant magnetic perturbations. *Phys. Plasmas*, 23(5):055703, 2016.
- [10] W. A. Newcomb. Magnetic differential equations. *Phys. Fluids*, 2(4):362, 1959.

- [11] S. R. Hudson, R. L. Dewar, G. Dennis, M. J. Hole, M. McGann, G. von Nessi, and S. Lazerson. Computation of multi-region relaxed magnetohydrodynamic equilibria. *Phys. Plasmas*, 19:112502, 2012.
- [12] L. Zakharov. Implementation of hamada principle in calculations of nested 3-d equilibria. *J. Plasma Phys.*, 81(6):515810609, 2015.

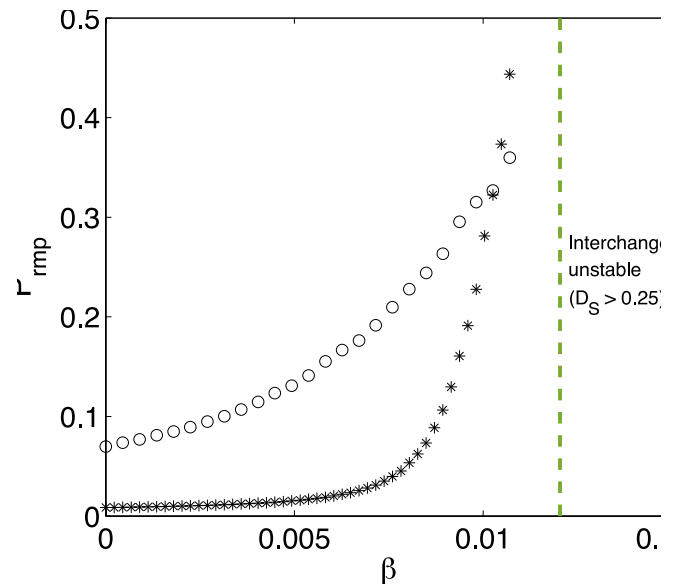


FIG. 3: Penetration of the perturbation on the resonant surface as a function of  $\beta$ , for  $\Delta t = 0.005$  (circles) and  $\Delta t = 0.001$  (stars). The vertical dashed line indicates where the Suydam condition is not satisfied. (Reproduced from [9].)

- [13] H. Weitzner. Ideal magnetohydrodynamic equilibrium in a non-symmetric topological torus. *Phys. Plasmas*, 21(2):022515, 2014.
- [14] S. R. Hudson, D. A. Monticello, A. H. Reiman, A. H. Boozer, D. J. Strickler, S. P. Hirshman, and M. C. Zarnstorff. Eliminating islands in high-pressure free-boundary stellarator magnetohydrodynamic equilibrium solutions. *Phys. Rev. Lett.*, 89(27):275003–1, 2002.
- [15] M. N Rosenbluth, R. Y. Dagazian, and P. H Rutherford. Nonlinear properties of the internal  $m=1$  kink instability in the cylindrical tokamak. *Phys. Fluids*, 16(11):1894, 1973.
- [16] R. L. Dewar, S. R. Hudson, A. Bhattacharjee, and Z. Yoshida. *Phys. Plasmas*, submitted, 2016.
- [17] J Loizu, S. R. Hudson, A. Bhattacharjee, and P. Helander. Magnetic islands and singular currents at rational surfaces in three-dimensional mhd equilibria. *Phys. Plasmas*, 22:022501, 2015.
- [18] W. A. Newcomb. Hydromagnetic stability of a diffuse linear pinch. *Ann. Phys.*, 10:232, 1960.
- [19] S. P. Hirshman and J. P Whitson. Steepest-descent moment method for three-dimensional magnetohydrodynamic equilibria. *Phys. Fluids*, 26(12):3553, 1983.
- [20] M. Taylor. A high performance spectral code for nonlinear MHD stability. *J. Comp. Phys.*, 110(2):407, 1994.
- [21] S. Lazerson, J Loizu, S.P. Hirshman, and S. R. Hudson. Verification of the ideal magnetohydrodynamic response at rational surfaces in the vmec code. *Phys. Plasmas*,

23:012507, 2016.

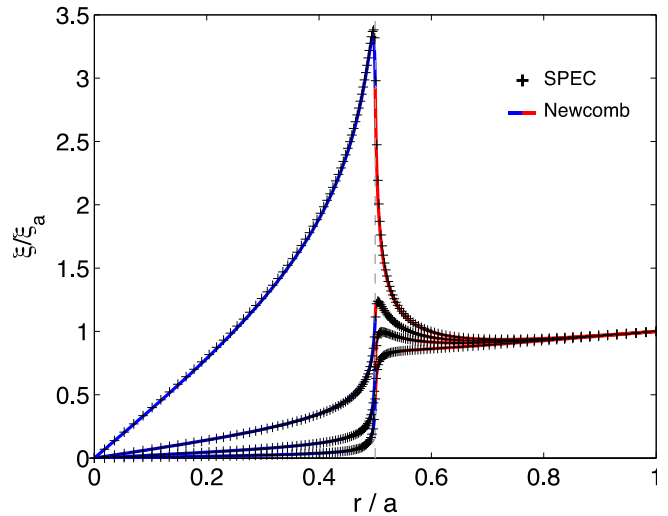


FIG. 4: SPEC linear solutions (crosses) and Newcomb solutions (solid lines) for  $\beta = 0\%$  (lower curve) to  $\beta = 1.1\%$  (upper curve), and for  $\Delta t = 0.0014$ . (Reproduced from [9].)